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$e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}(H_1)$ in the MSSM with explicit CP violations

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Abstract

We consider the effects of the CP-violating phases, *e.g.* $\arg(A_t)$ and $\arg(\mu)$ on the $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}(H_1)$ processes. The third generation squark trilinear terms give significant contributions to the Higgs potential at the one-loop level. This results in the changes of the stop masses, the Higgs mass and the lighter stop-the lighter stop*-the lightest Higgs coupling. We show the coupling and the loop effects on the processes. And we will discuss the determination method of soft parameters.

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I. INTRODUCTION

The Standard Model (SM) can explain the CP violation of the K system by means of the single KM phase in the CKM mixing matrix. And this paradigm will be tested in detail at B factories by measuring various CP violations in the B system. The latter will constitute an important test for the SM in the CP violating sector, and may give a hint for new physics related with CP violations beyond the KM phase [1] (which is strongly motivated in order to properly understand baryogenesis). The best motivated candidates for new physics are models with supersymmetry (SUSY), since SUSY can eliminate the quadratic divergences of the Higgs mass under grand unification theory (GUT). This is one of fine tuning problems of the Standard Model, the so–called gauge hierarchy problem. Therefore it would be natural to consider the effects of CP violation in the context of the supersymmetric theories. The simplest extension of the Standard Model is the Minimal Supersymmetric Standard Model (MSSM), which we consider in the following. The CP-violating parameters in the MSSM contribute to the electric dipole moments (EDMs) of electron and neutron, or ϵ_K depending on their flavor structures, so that these CP violating observables do constrain the CP violating phases in the MSSM. For example, the one loop contribution in the SUSY models to the neutron EDM is schematically given by $d_n \sim 2(100\,\mathrm{GeV}/\tilde{m})^2 \sin\phi \times 10^{-23} e\,\mathrm{cm}$, where $\tilde{m}$ is the overall SUSY scale and ϕ is the typical CP violating phase [2]. The latest experimental bound is $|d_n| < 6.3 \times 10^{-26} e\,\mathrm{cm}$ (90% C.L.) [3]. Therefore one can imagine various scenarios satisfying this tight experimental constraint. In the usual scenario, one assumes that the CP phase $|\phi| \leq 10^{-2}$ and the typical supersymmetric mass $\tilde{m} \leq 1$ TeV. Or ϕ can be $O(1)$ and the scalar masses of the first two generations are $O(1)$ TeV [1,2]— this is called the effective SUSY model. Recently, it was known that two independent CP violating phases in the minimal supergravity GUT model [4,5] can be $O(1 - 10^{-1})$ with the typical supersymmetric masses $\tilde{m} \leq 1$ TeV, if some internal cancellations among various contributions occur [6]. In this context of the internal cancellation mechanism of EDMs, two people [7] analyzed the CP violating phase effects on $gg \to \Phi^0$ ($\Phi^0 = h^0, H^0, A^0$) for the parameter space where the Higgs mixing effects [8] are negligible.

There are several important effects of the CP violating phases in the MSSM that have been studied recently. Firstly, the Higgs potential can be CP violating due to the one-loop effects of the third generation trilinear soft terms, although the tree level potential of the Higgs fields is CP invariant. So the lightest neutral Higgs mass can be larger or smaller than that of the CP conserving case (this can be phenomenologically important) [8], and the branching ratios of Higgs bosons can be changed significantly [9]. Secondly, the large CP violating phases may be coincident with the cosmological upper bound of the relic density of neutralinos [5], which are a candidate of dark matters in the R-parity invariant model. Thirdly, the chargino pair-production has a strong dependence of $\arg(\mu)$, but it still can be used in determination of some SUSY parameters such as $\tan\beta$, gaugino mass M_2, $|\mu|$ and $\cos(\arg(\mu))$ [10]. Fourthly, the direct CP asymmetry $A_{\rm CP}^{b\to s\gamma}$ of $B \to X_s\gamma$ can be as large as $\pm 16\%$ [11] (which is much larger than the SM contribution $\sim 0.5\%$ [12]), if chargino and stops are light enough. Therefore this may be a good place to look for a new source of CP violation. Also the CP phases contribute to the CP violating ϵ_K parameter of the neutral Kaon [2].

In the MSSM, the lighter stop may be the lightest squark [13], because the largeness of the

top Yukawa coupling makes (1) the diagonal components of stop mass^2 matrix smaller than those of two other generation squarks via renormalization group equations (RGEs) and (2) the off-diagonal components of stop mass^2 matrix larger than those of other squarks [14,15]. Therefore, lighter stop may be produced relatively easily in the next generation colliders. Furthermore, the stop has a large Yukawa coupling. This means that the associated Higgs productions with the stop-stop* pair could be sizable in the colliders [16]. These subjects have been considered by several groups in the MSSM without CP violations from the SUSY sector. In the presence of CP violating phases in A_t and μ parameters, the masses and the mixing parameters of stops and Higgs bosons, and also the stop-stop*-Higgs couplings would be modified. In this work, we consider these effects on the stop pair productions (with the lightest Higgs boson) at linear colliders.

This paper is organized as follows. In Sec. II, we study the one-loop effective potential of Higgs fields, the CP violating vacuum, the mixing of the different CP-property Higgs particles, the stop mass^2 matrix, and the stop-stop*-Higgs vertex. In Sec. III, we will choose an appropriate MSSM scenario satisfying the electron/neutron EDM constraints. And we analyze the CP phase dependences and the magnitudes of the relative vacuum angle ξ, stop mass and the cross sections of the $e^+e^- \rightarrow \tilde{t}_1\tilde{t}_1^{\,*}(H_1)$ processes. And we will discuss the determining method of the soft parameters in Sec. IV, and we conclude in Sec. V.

II. THE HIGGS AND STOPS SECTOR IN THE MSSM

A. Higgs sector

In the MSSM, it is well known that the Higgs F/D-term potentials are CP conserving at the tree level, and CP violating terms can reside in the soft SUSY breaking sector only. However, the squark-antisquark-Higgs trilinear soft terms can break CP, and they can generate *effective* CP violating terms in the Higgs effective potential via the squark loop corrections: namely, CP can be broken effectively in the Higgs sector. In terms of the Wilsonian action, the CP violating effective interactions occur when the higher frequency modes are integrated out in the CP-violating MSSM [17]. But the Wilsonian action may not be good, since it can be unphysical when a massless particle participates [18]. In Ref. [18], the electric charge of the scalar field in the Wilsonian action is dependent on the gauge fixing parameter ζ at the two-loop level in the massless scalar QED. This is related to the infrared effect. And the field-integrating-out effective action (*e.g.* the four Fermi interaction with W boson integrated out) is not good for this case either, since the integrated-out field, the lighter stop can be lighter than the heaviest Higgs field which can be included in the effective theory. Therefore, the 1PI effective action can be more suitable. The $\mathrm{U}(1)_Y$ toy model [19] can illustrate how the CP violating effective terms can be brought about by means of the 1PI effective action.

The 1-loop corrected effective potential of the Higgs fields (we will follow the notations of Ref. [8], where many important calculations were done) is given by

$$\begin{aligned}\mathcal{V}_{Higgs}^{eff} = {}& \mu_1^2\Phi_1^\dagger\Phi_1 + \mu_2^2\Phi_2^\dagger\Phi_2 + (m_{12}^2\Phi_1^\dagger\Phi_2 + h.c.) \\ & +\lambda_1(\Phi_1^\dagger\Phi_1)^2 + \lambda_2(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) \\ & +\lambda_5(\Phi_1^\dagger\Phi_2)^2 + \lambda_5^*(\Phi_2^\dagger\Phi_1)^2 + \lambda_6(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2) + \lambda_6^*(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_1)\end{aligned}$$

$$+\lambda_7(\Phi_2^\dagger\Phi_2)(\Phi_1^\dagger\Phi_2) + \lambda_7^*(\Phi_2^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) \ , \tag{1}$$

where $\lambda_i = 0$ $(i = 5, 6, 7)$ at the tree level so that they are generated entirely by quantum corrections, and Φ_i $(i = 1, 2)$ are the scalar components of the Higgs superfields. The $\Phi_2(\Phi_1)$ gives masses to the up-type (down-type) fermions. For $\tan\beta \sim O(1)$,

$$\begin{aligned} m_{12}^2 &= B\mu + l \cdot h_t^2 A_t \mu \ , \\ l &= \frac{1}{4\cdot 16\pi^2}\left(\frac{m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2}\ln\frac{m_{\tilde{t}_1}^2}{m_{\tilde{t}_2}^2} - 2\right) \ , \end{aligned} \tag{2}$$

where B is the soft parameter of the bilinear term, $h_t = \sqrt{2}\, m_t(\bar{m}_t)/v\sin\beta$ and $m_{\tilde{t}_i}$ $(i = 1, 2)$ are the masses of the lighter and heavier stops. The real parameter l is $O(10^{-3})$ for appropriate parameter ranges, and $B\mu$ is set to be real by the field-redefinition or $\mathrm{U}(1)_Q$ spurion transformation (this is our convention) [4]. Because the quantum correction is proportional to (Yukawa coupling)^2, the contribution of the 3rd generation squarks are larger than those of 1st/2nd generation squarks. Therefore, only stop contributions were written in Eq. (2) at low $\tan\beta$, for which the sbottom contributions become negligible. The parameter $B\mu$ denotes the mixing of two would-be CP-even Higgs bosons (scalar-scalar mixing), and $l \cdot h_t^2 A_t\mu$ represents the mixing of three CP-even and CP-odd Higgs bosons (scalar-pseudoscalar mixing). So m_{12}^2 plays an important role in the mixing of Higgs fields. If $|l \cdot h_t^2 A_t \mu| \ll |B\mu|$ and $\mu_1^2 \sim \mu_2^2 \sim |B\mu|$, we can expect that the scalar-pseudoscalar mixing is much smaller than the scalar-scalar mixing. Therefore, the mixing matrix $\mathcal{O}$ in Eq. (5) has a weak dependence of CP violating phases. In order to see the scalar-pseudoscalar mixing more clearly, the authors of Ref. [8] used $|A_t|$, $|\mu| \geq 1$ TeV, which may not be natural in the MSSM.

The $\mathrm{SU}(2)_L\times$ $\mathrm{U}(1)_Y$ gauge symmetry is broken spontaneously into $\mathrm{U}(1)_{\mathrm{em}}$ by the vacuum expectation values of two Higgs doublet fields [8]:

$$\begin{aligned} \Phi_1 &= \begin{pmatrix} \phi_1^+ \\ (v_1 + \phi_1 + i a_1)/\sqrt{2} \end{pmatrix} \ , \\ \Phi_2 &= e^{i\xi}\begin{pmatrix} \phi_2^+ \\ (v_2 + \phi_2 + i a_2)/\sqrt{2} \end{pmatrix} \ , \end{aligned} \tag{3}$$

where the VEVs v_i are real, and $\langle\Omega|\{\phi_i^+,\ \phi_i,\ a_i\}|\Omega\rangle = 0$ $(i = 1, 2)$. The relative phase ξ is determined from the minimum energy conditions of the Higgs potential [8], which are nothing but $T_\phi = \partial\mathcal{V}_{Higgs}^{eff}/\partial\phi = 0$ at the vacuum $|\Omega\rangle$ (ϕ represents a scalar field).

Due to the CP violating terms, *e.g.* $m_{12}^2\Phi_1^\dagger\Phi_2 + h.c.$, there exists a scalar-pseudoscalar transition, which results in the mixing of the Higgs fields with different CP quantum numbers in the limit of CP invariant Higgs sector. So the mass^2 matrix of the neutral Higgs particles [8] is

$$\mathcal{M}_H^2 = \begin{pmatrix} \mathcal{M}_P^2 & \mathcal{M}_{PS}^2 \\ \mathcal{M}_{SP}^2 & \mathcal{M}_S^2 \end{pmatrix} \ , \tag{4}$$

where $\mathcal{M}_{SP}^2 = {\mathcal{M}_{PS}^2}^T \neq 0$. Because the symmetries are spontaneously broken, there exists a *would-be* Goldstone boson, G^0. From the Goldstone's theorem, G^0 should be massless

for all the orders (from the tree level to all the loop levels). In the weak basis (G^0, $a = -a_1 \sin\beta + a_2 \cos\beta$, ϕ_1, ϕ_2), $\mathcal{M}^2_{H1j} = 0$ ($j = 2, 3, 4$) [8]. We will define a new matrix $\mathcal{M}^2_N$ as $\mathcal{M}^2_{Nij} = \mathcal{M}^2_{Hi+1,\,j+1}$ ($i, j = 1, 2, 3$). $\mathcal{M}^2_N$ is a real and symmetric matrix, so there exists a 3×3 orthogonal matrix $\mathcal{O}$ [8], which satisfies

$$\mathcal{O}^T \mathcal{M}^2_N \mathcal{O} = \mathrm{diag}(M^2_{H_3},\, M^2_{H_2},\, M^2_{H_1}) \ , \tag{5}$$

where $M_{H_3} \geq M_{H_2} \geq M_{H_1}$. The corresponding mass eigenstates, H_i ($i = 1, 2, 3$) , are

$$\begin{pmatrix} H_3 \\ H_2 \\ H_1 \end{pmatrix} = \mathcal{O}^T \begin{pmatrix} a \\ \phi_1 \\ \phi_2 \end{pmatrix} \ . \tag{6}$$

B. Stop sector

If we substitute the Φ_i's of Eqs. (3) into the MSSM Lagrangian, the stop mass^2 matrix $\mathcal{M}^2_{\tilde{t}}$ is in the

$$\begin{aligned} \mathcal{L}^{eff}_{mass} &= -\,(\tilde{t}^*_L\ \tilde{t}^*_R)\mathcal{M}^2_{\tilde{t}} \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix} \\ &= -\,(\tilde{t}^*_L\ \tilde{t}^*_R) \begin{pmatrix} m^2_{\tilde{t}L} & m^2_{\tilde{t}LR} \\ m^{2*}_{\tilde{t}LR} & m^2_{\tilde{t}R} \end{pmatrix} \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix} \ , \end{aligned} \tag{7}$$

where

$$\begin{aligned} m^2_{\tilde{t}L} &= M^2_{\tilde{t}_L} + m^2_t + M^2_Z \cos 2\beta \left(\frac{1}{2} - \frac{2}{3}\sin^2\theta_W \right) \ , \\ m^2_{\tilde{t}R} &= M^2_{\tilde{t}_R} + m^2_t + M^2_Z \cos 2\beta \cdot \frac{2}{3}\sin^2\theta_W \ , \\ m^2_{\tilde{t}LR} &= m_t \left(A^*_t e^{-i\xi} - \mu \cot\beta \right) \ . \end{aligned} \tag{8}$$

Although ξ is 0 or π (mod 2π) by the field redefinition at the tree level, it cannot be no longer the case when one includes the 1-loop effects [20], which is determined from the minimum energy conditions $T_\phi = \partial \mathcal{V}^{eff}_{Higgs}/\partial\phi = 0$ [8]. Therefore, the stop mixing angle $\theta_{\tilde{t}}$ [6] is changed via

$$\tan 2\theta_{\tilde{t}} = \frac{2|m^2_{\tilde{t}LR}|}{m^2_{\tilde{t}L} - m^2_{\tilde{t}R}} \ . \tag{9}$$

The relations between the mass and the weak eigenstates of stops are given by

$$\begin{aligned} \tilde{t}_1 &= \tilde{t}_L \cos\theta_{\tilde{t}} + \tilde{t}_R\, e^{-i\beta_{\tilde{t}}} \sin\theta_{\tilde{t}} \ , \\ \tilde{t}_2 &= -\tilde{t}_L\, e^{i\beta_{\tilde{t}}} \sin\theta_{\tilde{t}} + \tilde{t}_R \cos\theta_{\tilde{t}} \ , \end{aligned} \tag{10}$$

where $\beta_{\tilde{t}} = -\arg(m^2_{\tilde{t}LR})$. The mass eigenvalues of two stops are

$$m^2_{\tilde{t}_1,\tilde{t}_2} = \frac{m^2_{\tilde{t}L} + m^2_{\tilde{t}R} \mp \sqrt{(m^2_{\tilde{t}L} - m^2_{\tilde{t}R})^2 + 4|m^2_{\tilde{t}LR}|^2}}{2} \ . \tag{11}$$

Note that $m^2_{\tilde{t}_1,\tilde{t}_2}$ is dependent on the CP violating phases, $\arg(A_t)$ and $\arg(\mu)$ due to $m^2_{\tilde{t}LR}$ in Eqs. (8).

C. Stop-stop*-the lightest Higgs vertex

In the presence of CP violations in the Higgs sector, the stop-stop*-the lightest Higgs interaction is also modified. Defining the relevant interaction Lagrangian as

$$\mathcal{L}^{eff}_{\tilde{t}_1\tilde{t}_1{}^* H_1} = -V^{eff}_{\tilde{t}_1\tilde{t}_1{}^* H_1} \cdot \tilde{t}_1\tilde{t}_1{}^* H_1 \ , \tag{12}$$

the coupling $V^{eff}_{\tilde{t}_1\tilde{t}_1{}^* H_1}$ is given by

$$\begin{aligned} V^{eff}_{\tilde{t}_1\tilde{t}_1{}^* H_1} = & -\frac{g_W M_W}{2\cos^2\theta_W}(\mathcal{O}_{33}\sin\beta - \mathcal{O}_{23}\cos\beta)\left[\left(\frac{1}{2}-\frac{2}{3}\sin^2\theta_W\right)\cos^2\theta_{\tilde{t}} + \frac{2}{3}\sin^2\theta_W\sin^2\theta_{\tilde{t}}\right] \\ & + \frac{g_W m_t}{2M_W\sin\beta}\mathbf{Re}\left[A_t e^{i(\xi-\beta_{\tilde{t}})}(\mathcal{O}_{33}+i\mathcal{O}_{13}\cos\beta) - \mu\, e^{i\beta_{\tilde{t}}}(\mathcal{O}_{23}+i\mathcal{O}_{13}\cos\beta)\right] \\ & + \frac{g_W m_t^2}{M_W\sin\beta}\mathcal{O}_{33} \ . \end{aligned} \tag{13}$$

The g_W is the $\mathrm{SU}(2)_L$ gauge coupling constant and the $\mathcal{O}_{i3}$ are determined from $\mathcal{M}_N^2$ [8]. Since $\mathcal{M}_N^2$ depends on the CP violating phases, so do the $\mathcal{O}_{i3}$'s ($i = 1, 2, 3$). Note that at the tree level (without CP violations in the Higgs sector), $H_1 = h$, $H_2 = H$, $\xi = 0$ and the scalar-pseudoscalar mixing does not exist. Therefore, we recover the usual expressions,

$$\begin{aligned} \mathcal{O}_{13} &= 0 \ , \\ \mathcal{O}_{23} &= -\sin\alpha \ , \\ \mathcal{O}_{33} &= +\cos\alpha \ , \end{aligned}$$

where α is the scalar Higgs mixing angle at the tree level. Due to the 1-loop corrections, the tree level parameters are deformed, even inducing the mixing among Higgs bosons with different CP properties.

III. THE ANALYSES OF $E^+E^- \to \tilde{T}_1\tilde{T}_1{}^*(H_1)$

A. EDM constraints

Before we discuss the CP violating phase dependences of the stop masses and their production cross sections at next linear colliders (NLC's), we have to consider the neutron/electron EDM constraints on the CP violating phases, $\arg(A_t)$ and $\arg(\mu)$. In the previous section, $\arg(\mu)$ was transferred into $\arg(A_t)$ in the form of $\arg(A_t) + \arg(\mu)$ from Eqs. (2) and (8). So we can take μ to be real, and vary $\arg(A_t)$ from 0 to 2π in the following. In the minimal supergravity model, the phases $\arg(A_{t0})$ and $\arg(\mu_0)$ can be $O(1-10^{-1})$ without conflict with neutron/electron EDM constraints [6]. The A_{t0} and μ_0 are the values of A_t and μ at the GUT scale. But these two phases are rather correlated, namely having different signs, and $|\arg(A_{t0}) + \arg(\mu_0)|$ is typically smaller than $O(\pi/9)$ for low $\tan\beta$ and the soft mass scale $\gtrsim 500$ GeV [6]. Of course, we should consider the renormalization group behaviors of $\arg(A_t)$ and $\arg(\mu)$ down to the electroweak (EW) scale. In the ordinary phase convention of the minimal supergravity GUT (only A and μ terms have CP violating

phases), if an appropriate universality at the GUT scale is assumed, $\mathbf{Re}(A_{t\mathrm{EW}}) > \mathbf{Re}(A_{t0})$ and $\mathbf{Im}(A_{t\mathrm{EW}}) < \mathbf{Im}(A_{t0})$ from the typical RGEs and the numerical values of Ref. [14], where $A_{t\mathrm{EW}}$ is the value of A_t at the EW scale. This means that the $\arg(A_{t\mathrm{EW}}) < \arg(A_{t0})$ (about ten times smaller), which is different from the μ case, that is, $\arg(\mu_{\mathrm{EW}}) = \arg(\mu_0)$ since $\mathbf{Re}(\mu)$ and $\mathbf{Im}(\mu)$ have the same RGE, $\dot{y}(t) = a(t)y(t)$ [14]. So we can conclude that $|\arg(A_t) + \arg(\mu)|$ decreases from the GUT scale to the EW scale via RGEs. Therefore the phase effects on stop pair productions at NLC's would not be that prominent in this case, and thus is not proper for our analyses. This leads us to consider other scenarios with large CP phases. We will choose a slightly non-universal scenario for the trilinear couplings A_f [8,21], the underlying assumptions of which are [21]

- $|\arg(\mu)| \lesssim 10^{-2}$ from cosmological reasonings
- $|A_e|, |A_{u,c}|, |A_{d,s}| \lesssim 10^{-3}|\mu|$
- $A_t = A_b = A_\tau$, where the only large CP-violating phase is contained
- gluino mass $m_{\tilde{g}} \gtrsim 400$ GeV

The first two assumptions are chosen in order to to suppress the one-loop EDM contributions owing to $d_f^{1-\mathrm{loop}} \propto \{\mathbf{Im}(\mu), \mathbf{Im}(A_f)\}$ for electron and light quarks [21]. The two-loop contributions can be much smaller than the experimental bounds if $|\mu|$, $\tilde{m}_2 \gtrsim 100$ GeV [21]. And the neutron EDM has a significant three-loop [21] contribution $\frac{e\Lambda_\chi}{4\pi} d^G(\Lambda_\chi)$ from Weinberg's three gluon operator $\mathcal{O}_{3-\mathrm{gluon}} = -\frac{1}{6} d^G f_{abc} G^a_{\mu\rho} G^{b\rho}_{\nu} G^c_{\lambda\sigma} \epsilon^{\mu\nu\lambda\sigma}$, which comes from two-loop diagrams with top, stop and gluino internal lines [22,23]. The Λ_χ is the chiral symmetry breaking scale. It has been shown that the three-loop contribution $\propto 1/m^3_{\tilde{g}}$ are much smaller than the experimental bound of neutron EDM if gluino mass $m_{\tilde{g}}$ is larger than 400 GeV [22,23]. So the 3-loop contribution to the EDM can be sufficiently smaller than the experimental bound by the fourth assumption. In Ref. [21], large $\tan\beta$ scenarios ($40 \lesssim \tan\beta \lesssim 60$) with $|\mu|, |A_t| > 500$ GeV, $M_a \leq 500$ GeV and large CP phase are excluded, but low $\tan\beta$ scenarios ($\tan\beta \lesssim 20$) can be possible. We will consider the parameter ranges, which are not ruled out by the EDM constraints, *i.e.* $\tan\beta = 2(\ll 20)$, $|\mu| = 500,\ 1000$ GeV and $B = 30$ GeV, for which the would-be CP-odd Higgs boson has its mass $M_a \approx \sqrt{|2\mathbf{Re}(m^2_{12} e^{i\xi})/\sin 2\beta|} \approx 194,\ 274$ GeV, evading the EDM constraints by the authors of Ref. [21]. Also we choose $|A_t| = |\mu \cot\beta|$ in order to maximize the effects of the A_t phase in the stop mass matrix.

B. Vacuum angle ξ

We analyze the relative phase ξ of two Higgs VEVs in Eqs. (3). The ξ is a solution of the condition of a vanishing non-trivial tadpole parameter (see Ref. [8])

$$T_a = \frac{\partial \mathcal{V}^{eff}_{Higgs}}{\partial a} = -v \Big[\mathbf{Im}(m^2_{12} e^{i\xi}) + \mathbf{Im}(\lambda_5 e^{2i\xi}) v^2 \sin\beta\cos\beta + \frac{1}{2}\mathbf{Im}(\lambda_6 e^{i\xi}) v^2 \cos^2\beta + \mathbf{Im}(\lambda_7 e^{i\xi}) v^2 \sin^2\beta \Big] = 0. \tag{14}$$

In Fig. 1, we show ξ as a function of $\arg(A_t)$ for $\mu = +500$ GeV. ξ has two sets of values around 0 and π (mod 2π). If ξ is neither 0 nor π (mod 2π), CP is spontaneously broken by the vacuum due to the quantum corrections to the Higgs potential coming from the stop loop. However the effect on ξ is not that prominent since $\arg(m_{12}^2) \approx 0$ and $|m_{12}^2| \gg |\lambda_{5-7} v^2|$ for our parameter region in Eq. (14). In Fig. 2, we show ξ for the same set of parameters except that the sign of μ takes both positive and negative. The maximal deviations of ξ from 2π is again less than $O(2\pi/120)$ (*i.e.* has a weak phase dependence) and varies with the sign of μ. We calculated the results for a real μ with two different signs (*i.e.* $\arg(\mu) = 0$ or π) in order to see the effects of changing $\arg(\mu)$, even though such a negative sign of μ can not satisfy the first assumption of the slightly non-universal scenario, strictly speaking. If the magnitude of μ gets larger (about 1 TeV), the ξ has a smaller range of $\arg(A_t)$ as shown in Fig. 3, since the lightest Higgs mass M_{H_1} gets smaller and finally becomes imaginary elsewhere. This latter phenomenon was observed in Figs. 2(a) and 6(a) of Ref. [8]. Also the sizes of $\xi - \pi$(or 2π) become larger than those of the smaller μ case, as shown in Fig. 3.

In addition, ξ appears in the mass matrices of charginos and neutralinos. (see Refs. [14,24].) But the sizes of $\xi - \pi$(and 2π) are very small, so they can not have significant 1-loop effects on charginos and neutralinos. ξ may also appear in the strong CP parameter $\bar{\theta} = \theta_{\rm QCD} + \arg(\det(M_{\rm quark})) = \theta_{\rm QCD} + \arg(\det(M_u M_d))$ [25] through the Standard Model Yukawa terms since the masses of the up-type and down-type quarks are proportional to the VEVs of Higgs fields in Eqs. (3). The $M_{\rm quark}$ is an $N_f \times N_f$ matrix, where N_f is the number of flavors, and equal to 6 in the SM.

C. $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}$ process

Let us first consider a process, $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}$ at NLC's with $\sqrt{s} = 500$ GeV. As before, we take $\mu = 500$ GeV, $\tan\beta = 2$, $M_{\rm SUSY} = 500$ GeV, and finally $|A_t| = |\mu\cot\beta|$ in order to maximize the effects of $\arg(A_t)$. Otherwise, the effects of $\arg(A_t)$ relative to $\arg\mu$ will be negligible: see Eqs. (8), (9) and (11). Since $\sqrt{s} \gg M_Z$ at NLC's, we neglect M_Z in the Z boson propagator. The amplitude for this process with a left-handed initial electron is

$$\mathcal{M}^L_{\tilde{t}_1\tilde{t}_1^{\,*}} = \frac{g_W^2 \sin^2\theta_W}{s}(V_1 + V_2)\,\bar{v}_L(p')\,\gamma^\mu\,u_L(p)(k'_\mu - k_\mu)\ , \tag{15}$$

where p and p' are 4-momenta of e^- and e^+ respectively, $s = (p + p')^2$ and

$$V_1 = \frac{2}{3} - \frac{1}{\sin^2\theta_W\cos^2\theta_W}\left(\sin^2\theta_W - \frac{1}{4}\right)\left(\frac{1}{2}\cos^2\theta_{\tilde{t}} - \frac{2}{3}\sin^2\theta_W\right)\ , \tag{16}$$

$$V_2 = \frac{1}{4\sin^2\theta_W\cos^2\theta_W}\left(\frac{1}{2}\cos^2\theta_{\tilde{t}} - \frac{2}{3}\sin^2\theta_W\right)\ . \tag{17}$$

The amplitude with a right-handed electron $\mathcal{M}^R_{\tilde{t}_1\tilde{t}_1^{\,*}}$ is obtained by replacement $u_L,\ v_L \to u_R,\ v_R$ and $V_1 + V_2 \to V_1 - V_2$ in Eq. (15). We can measure the $\cos^2\theta_{\tilde{t}}$ from σ_L/σ_R [26], where σ_i $(i = L, R)$ are the cross sections corresponding to the polarizations of the electrons. In Ref. [26], this process was considered at the tree level in the MSSM with real μ and A_t. In the presence of A_t phase, the cross sections could be altered significantly if $|A_t| \sim |\mu\cot\beta|$, since the physical stop mass and the mixing angle $\cos^2\theta_{\tilde{t}}$ (V_1 and V_2) depend strongly on

the phase $\arg(A_t)$ through Eqs. (9) and (11) due to the fact that $|m^2_{\tilde{t}LR}|$ in Eqs. (8) can vary from 0 to $2m_t|A_t|$. Because the phase dependence of $m_{\tilde{t}_1}$ becomes prominent for our parameter region from Fig. 4, the cross sections can depend on the phase strongly, since $(1-4m^2_{\tilde{t}_1}/s)^{3/2}$ of the cross sections has a much stronger phase dependence through the stop mass than the coupling $(V_1^2+V_2^2)$. However, the 1-loop correction to the Higgs potential does not change the cross sections very much, which is not hard to understand. The phase dependences come from only A_t^* and $e^{i\xi}$ in Eq. (8). The 1-loop corrections hardly change ξ which has a very small amplitude, so the unique phase ξ that was generated by quantum corrections can be safely neglected. And, there is no significant change of the cross section as shown in Fig. 5.

The total cross section for $e^+e^- \to \tilde{t}_1\tilde{t}_1{}^*$ at $\sqrt{s}=500$ GeV is greater than 40 fb (see Fig. 5). Therefore, if a high integrated luminosity $\int \mathcal{L}dt = 500$ fb^{-1} could be achieved at $\sqrt{s}=500$ GeV as suggested by TESLA working group, we would expect more than about 20000 events for a year. This should be enough for studying the lightest stop mass, the cross section and the main decay modes of the lightest stop. The branching ratios of the stop in the CP conserving MSSM can be found in Ref. [27].

Since physical quantities such as stop masses and their production cross section depend strongly on $\arg(A_t)$, it is important to determine $\arg(A_t)$ experimentally. We will discuss this issue in Sec. IV.

D. $e^+e^- \to \tilde{t}_1\tilde{t}_1{}^*H_1$ process

We analyze the process $e^+e^- \to \tilde{t}_1\tilde{t}_1{}^*H_1$, where $\sqrt{s}=500$ GeV, $\mu=500,\ 1000$ GeV, $\tan\beta=2$, $M_{\rm SUSY}=500$ GeV and again $|A_t|=|\mu\cot\beta|$. At the tree level, the most significant diagrams have the $\tilde{t}_1\tilde{t}_1{}^*H_1$ vertex alone and other diagrams have $\tilde{t}_2\tilde{t}_1{}^{(*)}H_1$ and $Z^0\tilde{t}_1\tilde{t}_1{}^*$ vertices [26]. The 1-loop corrected amplitude with a left-handed electron is

$$\begin{aligned}\mathcal{M}^L_{\tilde{t}_1\tilde{t}_1{}^*H_1} = &-\frac{g_W^2\sin^2\theta_W\cdot V^{eff}_{\tilde{t}_1\tilde{t}_1{}^*H_1}}{s}\,(V_1+V_2)\,\bar{v}_L(p')\,\gamma^\mu\,u_L(p)\\ &\times\left(\frac{q'_\mu-k_\mu}{q'^2-m^2_{\tilde{t}_1}}+\frac{k'_\mu-q''_\mu}{q''^2-m^2_{\tilde{t}_1}}\right)\ ,\end{aligned}\tag{18}$$

where $q'=p+p'-k$, $q''=p+p'-k'$, k is the 4-momentum of $\tilde{t}_1$, and k' is the 4-momentum of $\tilde{t}_1{}^*$. The amplitude with a right-handed electron $\mathcal{M}^R_{\tilde{t}_1\tilde{t}_1{}^*H_1}$ is obtained by replacement $u_L,\ v_L \to u_R,\ v_R$ and $V_1+V_2 \to V_1-V_2$ in Eq. (18). The cross sections depend on the CP violating phase because $V^{eff}_{\tilde{t}_1\tilde{t}_1{}^*H_1}$ and $\cos^2\theta_{\tilde{t}}$ are varying with the phase via Eqs. (4), (5), (8) and (9).

As noticed in the previous subsection, we can neglect the 1-loop effect coming through ξ due to its tiny variations around the tree level values. However, the quantum corrections can affect the $\mathcal{O}_{i3}$ $(i=1,2,3)$ and the Higgs mass M_{H_1} significantly, in general. But in Sec. II, it was argued that the mixing matrix $\mathcal{O}$ can have a weak phase dependence in some region, and our parameters are exactly in that region. Therefore, only the lightest Higgs mass M_{H_1} is much affected by the quantum corrections (M_{H_1} is independent of $\arg(A_t)$ at the tree level). This results in a sizable change of the cross section for $e^+e^- \to \tilde{t}_1\tilde{t}_1{}^*H_1$ through the

strong dependence of the Higgs mass on the A_t phase. The numerical results are shown in Fig. 6 (in filled circles) along with the tree level results (in open circles). Also the possible range for the cross section becomes much wider once we include the one-loop corrections to the Higgs potential. The symmetry of the cross section about $\arg(A_t) = \pi$ in Fig. 6 is represented by the symmetry under $\arg(A_t) \to 2\pi - \arg(A_t)$, *i.e.* $\cos(\arg(A_t)) \to \cos(\arg(A_t))$ and $\sin(\arg(A_t)) \to -\sin(\arg(A_t))$. The Higgs masses [8] are

$$M^2_{H_i} = -\frac{1}{3} r + 2 \left(-\frac{p^3}{27} \right)^{1/6} \cos\left(\frac{\phi}{3} + \delta_i \right) \; , \tag{19}$$

where

$$\begin{aligned}
r &= -\mathrm{tr}(\mathcal{M}^2_N) \; , \\
s &= \left[\mathrm{tr}^2(\mathcal{M}^2_N) - \mathrm{tr}(\mathcal{M}^4_N) \right] /2 \; , \\
t &= -\det(\mathcal{M}^2_N) \; , \\
p &= (3s - r^2)/3 \; , \\
q &= 2r^3/27 - rs/3 + t \; , \\
\phi &= \cos^{-1} \left(q/2\sqrt{-p^3/27} \right) \; , \\
\delta_i &= 0, \pm 2\pi/3 \; (i = 1,2,3) \; .
\end{aligned} \tag{20}$$

Note that $\mathcal{M}^2_N$ is of form $f(|B|, \xi, |A_t|, \arg(A_t), |A_b|, \arg(A_b), M^2_{\mathrm{SUSY}})$ without μ^2_i $(i = 1, 2)$ by using two minimum energy conditions. Since the stop mass $m_{\tilde{t}_1}$ and the Higgs mass M_{H_1} are symmetric under $\arg(A_t) \to 2\pi - \arg(A_t)$ [8], the cross section has that symmetry. And the cross sections have a minimum near $\arg(A_t) = \pi$ in Fig. 6 because the coupling $(V^{eff}_{\tilde{t}_1 \tilde{t}_1^* H_1})^2 (V_1^2 + V_2^2)$ has a minimum at $\arg(A_t) = \pi$ and its phase dependence is stronger than the dependence through the masses of stop and Higgs boson unlike the $e^+e^- \to \tilde{t}_1 \tilde{t}_1^*$ process. For larger μ (about 1 TeV), the qualitative features remain the same as the lower μ case, as shown in Fig. 7. But the larger $|\mu|$ brings about the region of $\arg(A_t)$ where the lightest Higgs mass M_{H_1} becomes imaginary (*i.e.* unphysical) as discussed before.

And the 1-loop cross sections are sufficiently enhanced from the tree level results. This can be a good news in the light of experiment. For the integrated luminosity, $\int \mathcal{L} dt = 500$ fb^{-1} at $\sqrt{s} = 500$ GeV, one can have more than 250 events for a year, since the cross section is larger than 0.5 fb from Fig. 6. The study of the decay modes and the branching ratios of Higgs bosons in the presence of μ and A_t phases can be found in Ref. [9].

IV. THE DETERMINATIONS OF THE SOFT PARAMETERS

In the previous section, we observed important physical quantities (stop mass and cross sections) have strong phase dependences. Therefore, it is important to determine the value of the CP violating phase $\arg(A_t)$. In order to fix n unknown parameters, we need n independent equations involving these unknown parameters. The equations should be, of course, about physical observables like masses, cross sections and so on. Therefore, if we have enough independent physical quantities from experiments, the soft parameters may be

determined in principle. For example, it is well known that $\tan\beta$, $|\mu|$ and $\arg(\mu)$ can be determined from $e^+e^- \to \tilde{\chi}_i{}^+\tilde{\chi}_j{}^-$ with $i,j=1,2$ [10]. This was possible because charginos are spin-half particles, so that there are many physical observables with different combinations of chargino helicities [10]. Therefore, we will assume these 3 parameters are known in the following. Also $B\mu$ was set to be real by the field-redefinition (this is our convention) [4], and then the soft parameter $\arg(B)$ should be fixed from $\arg(\mu)$ which is determined from $e^+e^- \to \tilde{\chi}_i{}^+\tilde{\chi}_j{}^-$.

In the process $e^+e^- \to \tilde{t}_i\tilde{t}_j{}^*$ $(i,j=1,2)$, there are 5 unknown parameters, ξ, $M^2_{\tilde{t}_L}$, $M^2_{\tilde{t}_R}$, $|A_t|$ and $\arg(A_t)$. However, the outgoing particles in this case are scalar unlike the case of the chargino pair production. Therefore only three equations are possible from the measured values of $\cos^2\theta_{\tilde{t}}$, $m^2_{\tilde{t}_1}$ and $m^2_{\tilde{t}_2}$, which are functions of $\xi, M^2_{\tilde{t}_L}, M^2_{\tilde{t}_R}, |A_t|, \arg(A_t)$. (The value of $\cos^2\theta_{\tilde{t}}$ will be determined from σ_L/σ_R [26].) Therefore, it is impossible to find the parameters of the stop sector from this process alone. In addition, since ξ depends on the Higgs sector, the process $e^+e^- \to \tilde{t}_i\tilde{t}_j{}^*$ $(i,j=1,2)$ alone can not provide sufficient informations to fix these five unknown parameters. Other processes are needed. Anyway, we have three other independent equations, *i.e.* the four minimum energy conditions ($T_\phi = \partial\mathcal{V}^{eff}_{Higgs}/\partial\phi = 0$), one of which is trivial from the Goldstone's theorem [19,20], that is to say, there remain only three equations. Those equations have 9 real parameters, $|B|$, ξ, μ_1^2, μ_2^2, $|A_t|$, $\arg(A_t)$, $|A_b|$, $\arg(A_b)$ and $M^2_{\rm SUSY}$ (see Ref. [8]). For example, one equation (*c.f.* Eq. (14)) is

$$T_a = \frac{\partial\mathcal{V}^{eff}_{Higgs}}{\partial a} = 0 = -v\left[\mathbf{Im}(m^2_{12}e^{i\xi}) + \mathbf{Im}(\lambda_5 e^{2i\xi})v^2\sin\beta\cos\beta + \frac{1}{2}\mathbf{Im}(\lambda_6 e^{i\xi})v^2\cos^2\beta \right. \\ \left. +\mathbf{Im}(\lambda_7 e^{i\xi})v^2\sin^2\beta\right] , \tag{21}$$

where

$$\lambda_5 = \frac{h_t^4\mu^2A_t^2}{M^4_{\rm SUSY}} f_5(h_t,h_b,g_s,M^2_{\rm SUSY},\bar{m}_t{}^2) + \frac{h_b^4\mu^2A_b^2}{M^4_{\rm SUSY}} g_5(h_t,h_b,g_s,M^2_{\rm SUSY},\bar{m}_t{}^2) \tag{22}$$

$$\lambda_6 = \frac{h_t^4|\mu|^2\mu A_t^2}{M^4_{\rm SUSY}} f_6(h_t,h_b,g_s,M^2_{\rm SUSY},\bar{m}_t{}^2) \\ +\frac{h_b^4\mu}{M_{\rm SUSY}}\left(\frac{6A_b}{M_{\rm SUSY}} - \frac{|A_b|^2A_b}{M^3_{\rm SUSY}}\right) g_6(h_t,h_b,g_s,M^2_{\rm SUSY},\bar{m}_t{}^2) , \tag{23}$$

$$\lambda_7 = \frac{h_b^4|\mu|^2\mu A_b^2}{M^4_{\rm SUSY}} f_7(h_t,h_b,g_s,M^2_{\rm SUSY},\bar{m}_t{}^2) \\ +\frac{h_t^4\mu}{M_{\rm SUSY}}\left(\frac{6A_t}{M_{\rm SUSY}} - \frac{|A_t|^2A_t}{M^3_{\rm SUSY}}\right) g_7(h_t,h_b,g_s,M^2_{\rm SUSY},\bar{m}_t{}^2) , \tag{24}$$

where f_i and g_i $(i=5,6,7)$ are of form

$$f(x_1,x_2,x_3,x_4,x_5) = k\left[1 - \frac{1}{16\pi^2}\left(ax_1^2 + bx_2^2 + 16x_3^2\right)\ln\frac{x_4}{x_5}\right] \tag{25}$$

and k is $O(10^{-3})$ and a, b are $O(1)$. For a small mass^2 splitting between the stop mass eigenstates, $M^2_{\rm SUSY} = \frac{1}{2}(m^2_{\tilde{t}_1}+m^2_{\tilde{t}_2}) = \frac{1}{2}{\rm tr}(\mathcal{M}^2_{\tilde{t}})$ (see Ref. [28] and the Appendix of Ref. [8]), so $M^2_{\rm SUSY}$ is known. Three parameters ξ, $|A_t|$ and $\arg(A_t)$ were already counted when we considered the processes $e^+e^- \to \tilde{t}_i\tilde{t}_j{}^*$ $(i,j=1,2)$. Therefore, there are 11 real parameters

and 7 independent equations from the process $e^+e^- \to \tilde{t}_i\tilde{t}_j^{\,*}$ $(i,j=1,2)$ and the vanishing tadpole conditions. Still four parameters remain undetermined. However, for small $\tan\beta = O(1)$, the sbottom trilinear coupling can be ignored, *i.e.* $h_b \simeq 0$. As a result, the effects of $|A_b|$ and $\arg(A_b)$ can be effectively neglected. Thus only two parameters are left undetermined for low $\tan\beta$. At this point, let us note that the masses of two neutral Higgs bosons give two independent equations from Eq. (19). The masses M_{H_i} of the light Higgs bosons can be determined from $e^+e^- \to H_iZ^0$ $(i=1,2)$ [8,29–31] (since the H_1Z^0 production gives an information about the Higgs mass M_{H_1} and has a cross section at $\sqrt{s}=500$ GeV about $O(10)$ times larger than the $\tilde{t}_1\tilde{t}_1^{\,*}H_1$ production (for example, see Ref. [31]), the associated Higgs production may not give more information), $e^+e^- \to H_iH_j$ $(i,j=1,2)$ [32], $u\bar{d} \to H_iW^+$ $(i=1,2)$ [8,33]. Therefore, once the neutral Higgs masses are determined by the experiments, we will have two necessary informations, and finally we can determine all the parameters. For intermediate $\tan\beta \lesssim O(20)$, h_b can not be neglected and two parameters $|A_b|$ and $\arg A_b$ remain to be fixed. In such a case, the masses of two other charged and neutral Higgs bosons will provide the necessary informations.

Of course, we have to consider the lower bound of the lightest Higgs mass and the EDM constraints for the above arguments. As we have discussed above, the determinations are not very simple like the chargino pair production [10] since the equations are more complicated.

V. CONCLUSIONS

We have analyzed how the CP violating phase in A_t parameter can affect the relative phase ξ of two Higgs fields, the stop masses and cross sections in the two processes $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}(H_1)$, which may be within the scope of the next linear colliders. The relative phase (ξ) of the two Higgs fields that arise from the quantum corrections has very small dependence on $\arg(A_t)$ and $|\mu|$. The sizes of $\xi-\pi$(and 2π) are also numerically very small. In the process $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}$, the stop mixing angle $\theta_{\tilde{t}}$, the stop masses $m_{\tilde{t}_i}$ $(i=1,2)$ and the cross sections have very little dependence on ξ, but strong dependence on $\arg(A_t)$ since the stop mass itself changes a lot when $\arg(A_t)$ varies. The typical cross section is order of a several tens of fb, which is well within the scope of NLC's. In the process $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}H_1$, the loop correction becomes very important, since the lightest Higgs boson mass can be strongly affected by loop corrections with complex trilinear coupling A_t. However the neutral Higgs boson mixing matrix or ξ does not change significantly even if quantum corrections are included. In this case, the typical cross section is order of $O(0.1-10)$ fb depending on the A_t phase and $|\mu|$. Therefore this mode can be (less) promising at NLC's depending on the soft SUSY breaking parameters.

Since important physical quantities (stop mass and cross sections) may depend on the CP phases of soft SUSY breaking parameters, it is necessary to have a strategy for determination of these parameters in the presence of CP violating phase. We argued that, for low $\tan\beta$ region, it is in fact possible to fix A_t and other mass parameters in the Higgs and stop sectors through $e^+e^- \to \tilde{t}_i\tilde{t}_j^{\,*}$, the vanishing tadpole conditions and the masses of neutral Higgs bosons. For intermediate $\tan\beta$ upto ~ 20, the situation is the same if one has informations on the masses of two other charged and neutral Higgs bosons. The whole procedure to determine soft parameters seems complicated, but it will be possible.

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FIGURES

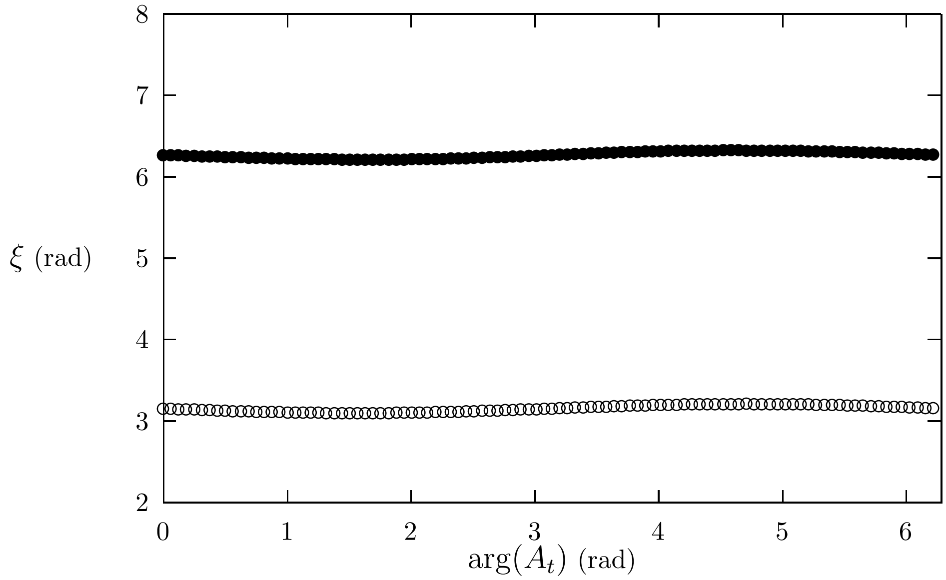

FIG. 1. The variations of ξ with respect to $\arg(A_t)$ ($\mu = 500$ GeV, $|A_t| = 250$ GeV, $M_a \approx 194$ GeV and $\tan\beta = 2$) The $\bullet$ denotes the point which corresponds to $\xi(\arg(A_t) = 0) = 2\pi$ (or 0) at the tree level and the $\circ$ denotes the point which corresponds to $\xi(\arg(A_t) = 0) = \pi$ at the tree level.

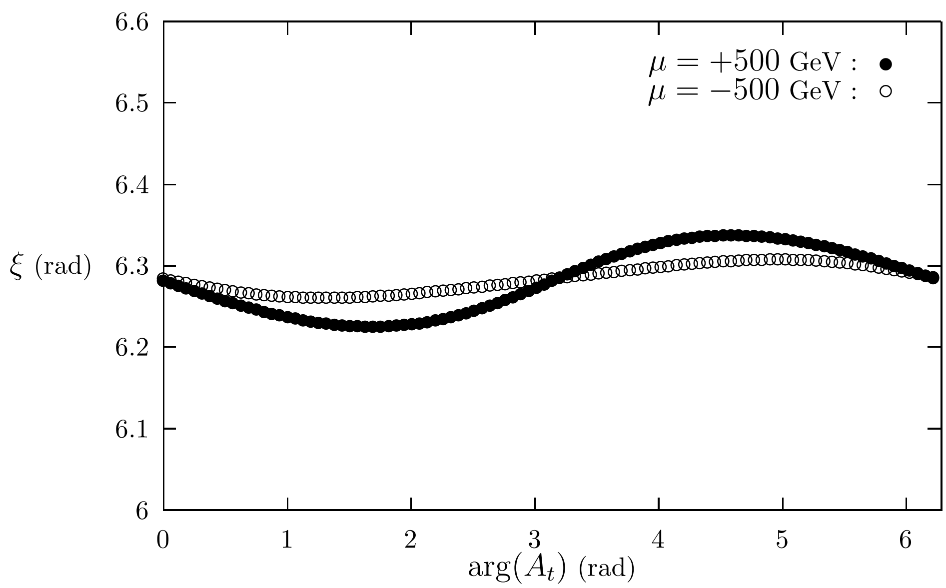

FIG. 2. The variations of ξ with respect to the sign of μ (or $\arg(\mu) = 0, \pi$) ($\mu = \pm 500$ GeV, $|A_t| = 250$ GeV, $\tan\beta = 2$ and $M_{\rm SUSY} = 500$ GeV)

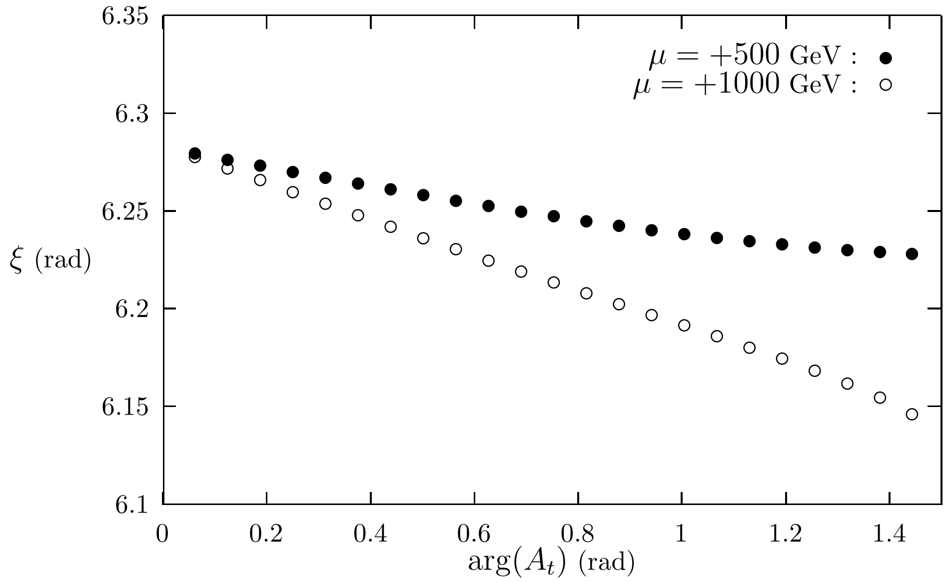

FIG. 3. The variations of ξ's with respect to the magnitude of μ ($|A_t| = |\mu \cot\beta|$, $\tan\beta = 2$ and $M_{\rm SUSY} = 500$ GeV) Note that the range of $\arg(A_t)$ is [0, 1.5] approximately, because the lightest Higgs mass M_{H_1} is imaginary elsewhere owing to the large $|\mu|$.

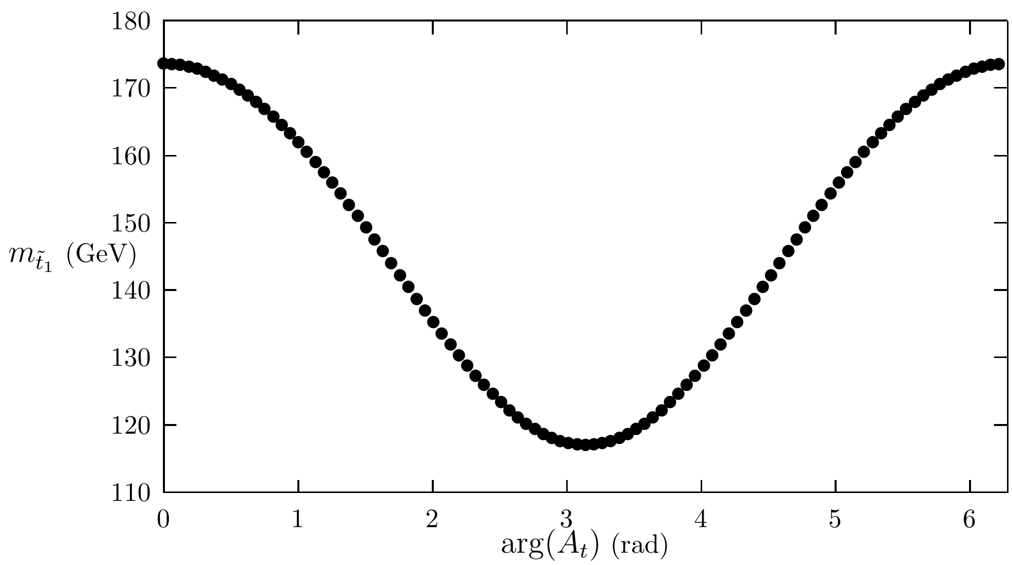

FIG. 4. The variations of the lightest stop mass with respect to $\arg(A_t)$ ($\mu = 500$ GeV, $|A_t| = 250$ GeV, $\tan\beta = 2$ and $M_{\rm SUSY} = 500$ GeV)

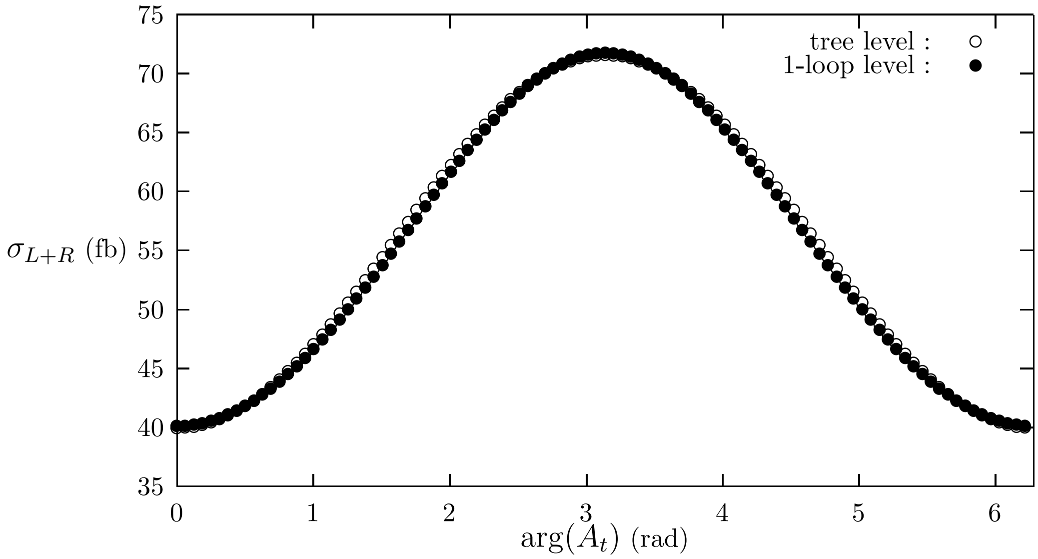

FIG. 5. The variations of the total cross sections with respect to the $\arg(A_t)$ in the process $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}$ ($\sqrt{s} = 500$ GeV, $\mu = 500$ GeV, $|A_t| = 250$ GeV, $M_a \approx 194$ GeV, $\tan\beta = 2$ and $M_{\rm SUSY} = 500$ GeV)

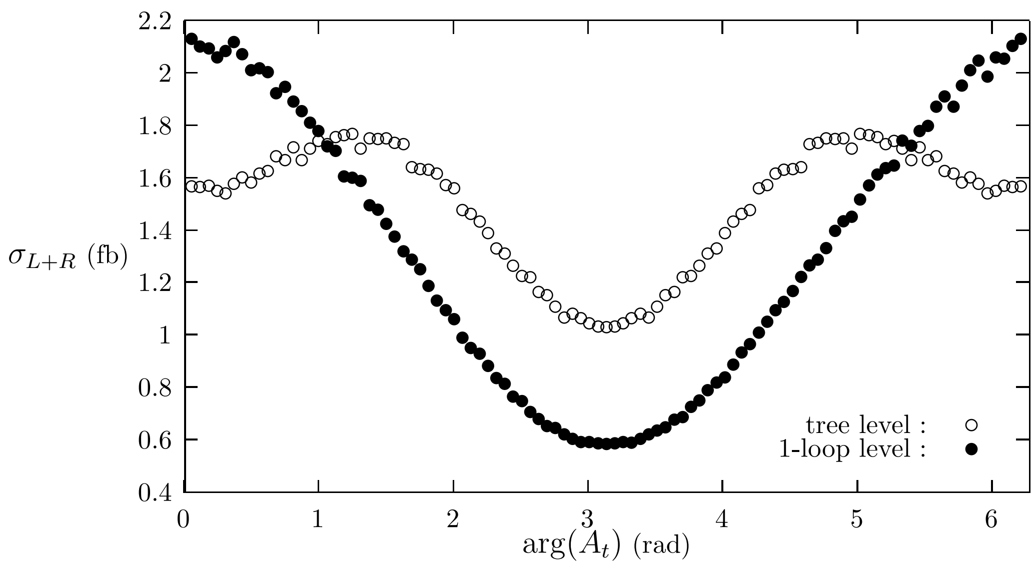

FIG. 6. The variations of the total cross sections with respect to $\arg(A_t)$ in the process $e^+e^- \to \tilde{t}_1\tilde{t}_1^{\,*}H_1$ ($\sqrt{s} = 500$ GeV, $\mu = 500$ GeV, $|A_t| = 250$ GeV, $M_a \approx 194$ GeV, $\tan\beta = 2$ and $M_{\rm SUSY} = 500$ GeV)

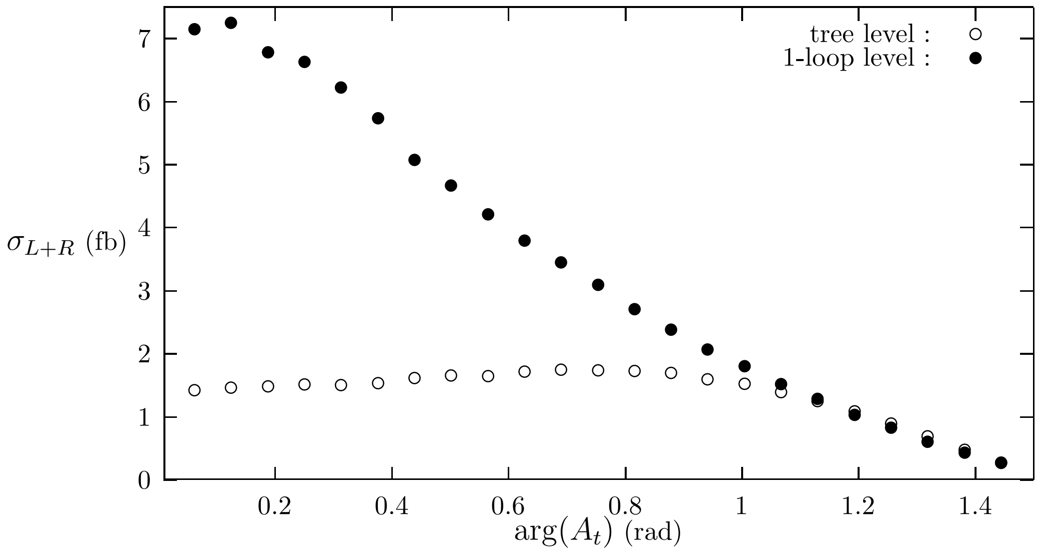

FIG. 7. The variations of the total cross sections with respect to $\arg(A_t)$ in the process $e^+e^- \to \tilde{t}_1\tilde{t}_1{}^*H_1$ ($\sqrt{s} = 500$ GeV, $\mu = 1000$ GeV, $|A_t| = 500$ GeV, $M_a \approx 274$ GeV, $\tan\beta = 2$ and $M_{\rm SUSY} = 500$ GeV) Note that the range of $\arg(A_t)$ is [0, 1.5] approximately, because the lightest Higgs mass M_{H_1} is imaginary elsewhere owing to the large $|\mu|$.